Quantum transport in mesoscopic systems and the measurement problem

S. A. Gurvitz

Department of Particle Physics, Weizmann Institute of Science, Rehovot 76100, Israel

(February 1, 2008)

Abstract

We study noninvasive measurement of stationary currents in mesoscopic systems. It is shown that the measurement process is fully described by the Schrödinger equation without any additional reduction postulate and without the introduction of an observer. Nevertheless the possibility of observing a particular state out of coherent superposition leads to collapse of the wave function, even though the measured system is not distorted by interaction with the detector. Experimental consequences are discussed.

PACS numbers: 03.65.Bz, 73.20.Dx, 73.23.Hk

According to the principles of quantum mechanics, a system in a linear superposition of several states undergoes a collapse to one of the states after measurement. More precisely, the density matrix of the measured system, $\rho(t)=\sum_{nm}|n\rangle\rho_{m,n}\langle m|$ collapses to the statistical mixture, $\rho(t)\to\sum_m|m\rangle\rho_{m,m}\langle m|$. This is the von Neumann projection postulate [1]. Since both the measuring device (the *detector*) and the measured system are described by the Schrödinger equation, the question arises of how such a non-unitary process takes place. The problem becomes even more acute when the measured system is a macroscopic one [2]. Then the distortion due to interaction with the detector can be made negligibly small — noninvasive measurement — so that the collapse mechanism appears even more mysterious [3].

A weak point of many studies of the measurement problem is the lack of a detailed quantum mechanical treatment of the entire system, that is, of the detector and the measured system together. The reason is that the detector is usually a macroscopic system, the quantum mechanical analysis of which is very complicated. Mesoscopic systems might thus be more useful for study of the measurement problem. In this paper we discuss the measurement of stationary processes — dc currents — in mesoscopic systems by using the recently derived quantum rate equations for quantum transport [4–6]. These equations applied to the entire system, allow us to follow the measurement process in great detail.

Let us consider quantum transport in small tunneling structures (quantum dots). These systems have attracted great attention due to the possibility of investigating single-electron effects in the electric current [7]. Until now research has been mostly concentrated on single dots, but the rapid progress in microfabrication technology has made it possible the extension to coupled dot systems with aligned levels [8,9]. In contrast with a single dot, the electron wave function inside a coupled dot structure is a superposition of electron states localized in each of the dots. The collapse of the wave function and its influence on the resonant current can thus be studied in these systems with a detector showing single electron charging of a quantum dot. Such a detector can be realized as a separate measuring circuit near the measured system [10,11].

A possible setup is shown schematically in Fig. 1. Two quantum dots, represented by quantum wells, are coupled to two separate reservoirs at zero temperature. The resonant levels E_0 and E_1 are below the corresponding Fermi levels. In the absence of electrostatic interaction between electrons, the dc resonant currents in the upper well (the detector) and the lower well (the measured system) are respectively

$$I_D^{(0)} = e\frac{\Gamma_0}{2}, \qquad\qquad I_S^{(0)} = e\frac{\Gamma_L\Gamma_R}{\Gamma_L+\Gamma_R} \tag{1}$$

The situation is different in the presence of electron-electron interaction between the dots, $H_{int} = Un_0n_1$, where $n_{0,1}$ are the occupancies of the upper and the lower wells and U is the Coulomb repulsion energy. If $E_0+U > \tilde{E}_F^L$, an electron from the left reservoir cannot enter the upper dot when the lower dot is occupied [Fig. 1 (c,d)]. On the other hand, the displacement of the level E_1 is much less important, since it remains below the Fermi level, $E_1+U < E_F^L$. The upper dot can thus be considered as a detector, registering electrons entering the lower dot [11]. For instance, by measuring the detector current, I_D, one can determine the current in the lower dot. The same setup can be used for measurement of the current in the coupled dot system shown in Fig. 2.

The quantum transport in the structures described above is fully determined by evolution of the density matrix for the entire system: $i\dot{\rho}(t) = [\mathcal{H},\rho]$, for $\mathcal{H} = H_D + H_S + H_{int}$, where $H_{D,S}$ are the tunneling Hamiltonians of the detector and the measured system, respectively, and $H_{int} = \sum_i U_i n_0 n_i$ describes their mutual Coulomb interaction. n_0 and n_i are the occupancies of the detector and the dot i of the multi-dot system. Using this equation one can find the current in the detector (or in the measured system) as the time derivative of the total average charge $Q(t)$ accumulated in the corresponding right reservoir (collector): $I(t) = \dot{Q}(t)$, where $Q(t) = e\mathrm{Tr}[\rho^R(t)]$, and $\rho^R(t)$ is the density matrix of the right reservoir. It was shown [4,6] that $I(t)$ is directly related to the density matrix of the multi-dot system $\sigma(t)$, obtained from the total density matrix $\rho(t)$ by tracing out the reservoir states. One finds that the current in the detector or in the measured system is given by

$$I(t) = e\sum_c \sigma_{cc}(t)\Gamma_R^{(c)}, \tag{2}$$

where $\sigma_{cc} \equiv \langle c|\sigma|c\rangle$ and the sum is over states $|c\rangle$ in which the well adjacent to the corresponding collector is occupied. $\Gamma_R^{(c)}$ is the partial width of the state $|c\rangle$ due to tunneling to the collector. The density matrix σ_{ij} obeys the following system of modified rate equations [6],

$$\dot{\sigma}_{aa} = i \sum_{b(\neq a)} \Omega_{ab}(\sigma_{ab} - \sigma_{ba}) - \sigma_{aa} \sum_{d(\neq a)} \Gamma_{a\to d} + \sum_{c(\neq a)} \sigma_{cc}\Gamma_{c\to a} \ , \tag{3a}$$

$$\begin{aligned} \dot{\sigma}_{ab} &= i(E_b - E_a)\sigma_{ab} + i\left(\sum_{b'(\neq b)} \sigma_{ab'}\Omega_{b'b} - \sum_{a'(\neq a)} \Omega_{aa'}\sigma_{a'b}\right) \\ &\quad -\frac{1}{2}\sigma_{ab}\left(\sum_{d(\neq a)} \Gamma_{a\to d} + \sum_{d(\neq b)} \Gamma_{b\to d}\right) + \frac{1}{2}\sum_{a'b'\neq ab} \sigma_{a'b'}\left(\Gamma_{a'\to a} + \Gamma_{b'\to b}\right) . \end{aligned} \tag{3b}$$

These equations were obtained from the many-body Schroödinger equation by integrating out the reservoir states, and assuming that the energy levels in the dots are not very close to the Fermi levels in the reservoirs. Here Ω_{ij} denote hopping matrix elements of the tunneling Hamiltonian and $\sigma_{ba} = \sigma^*_{ab}$. The width $\Gamma_{a\to b}$ is the probability per unit time for the multi-dot system to make a transition from the state $|a\rangle$ to the state $|b\rangle$ due to the tunneling to (or from) the reservoirs, to interaction with the phonon bath, or to any other interaction generated by a continuum-state medium. The non-diagonal matrix elements are determined by Eq. (3b), which resembles the optical Bloch equation supplemented with the last term. The latter contributes only for those *one-electron* transitions that convert at the same time the state $|a'\rangle$ into $|a\rangle$ *and* the state $|b'\rangle$ into $|b\rangle$.

It follows from Eq. (2) that the current is determined by the diagonal elements of the density matrix $\sigma(t)$. The non-diagonal elements ("coherences") influence the current via their coupling with the diagonal matrix elements (the first term in Eq. (3a)). This coupling is due to transitions among discrete states. In the absence of such transitions, for instance in the system shown in Fig. 1, the diagonal and non-diagonal matrix elements are decoupled in the rate equations. In this case the electron current is described by the classical rate equations. Note, this does *not* imply that the coherences vanish.

Let us apply Eqs. (3) to the system shown in Fig. 1. For simplicity we disregard the spin degree of freedom. There are four available states of the device: $|a\rangle$ – both wells are

empty; $|b\rangle$ – the upper well is occupied; $|c\rangle$ – the lower well is occupied; $|d\rangle$ – both wells are occupied. The rate equations for the diagonal matrix elements, obtained from Eq. (3a), are:

$$\dot{\sigma}_{aa} = -(\Gamma_0 + \Gamma_L)\sigma_{aa} + \Gamma_0\sigma_{bb} + \Gamma_R\sigma_{cc} \tag{4a}$$

$$\dot{\sigma}_{bb} = -(\Gamma_0 + \Gamma'_L)\sigma_{bb} + \Gamma_0\sigma_{aa} + \Gamma'_R\sigma_{dd} \tag{4b}$$

$$\dot{\sigma}_{cc} = -\Gamma_R\sigma_{cc} + \Gamma_L\sigma_{aa} + 2\Gamma'_0\sigma_{dd} \tag{4c}$$

$$\dot{\sigma}_{dd} = -(2\Gamma'_0 + \Gamma'_R)\sigma_{dd} + \Gamma'_L\sigma_{bb}, \tag{4d}$$

which are supplemented with the initial condition, $\sigma_{aa}(0) = 1$ and $\sigma_{bb}(0) = \sigma_{cc}(0) = \sigma_{dd}(0) = 0$.

The currents in the detector and in the lower well, Eq. (2), are respectively $I_D(t)/e = \Gamma_0\sigma_{bb}(t) + \Gamma'_0\sigma_{dd}(t)$ and $I_S(t)/e = \Gamma_R\sigma_{cc}(t) + \Gamma'_R\sigma_{dd}(t)$. The stationary (dc) current corresponds to $I = I(t \to \infty)$.

Solving Eqs. (4) we find in the limit $\Gamma_0, \Gamma'_0 \gg \Gamma_{L,R}, \Gamma'_{L,R}$

$$I_S/e = \frac{\Gamma_R(\Gamma_L + \Gamma'_L)}{\Gamma_L + \Gamma'_L + 2\Gamma_R}\,, \qquad\qquad \frac{I_D}{I_S} = \frac{\Gamma_0}{\Gamma_L + \Gamma'_L}\,, \tag{5}$$

The measurement of the detector current I_D can thus be considered as a measurement of the current in the lower well, I_S. However, I_S, Eq. (5), is distorted by the detector [$I_S \neq I_S^{(0)}$, Eq. (1)] since $\Gamma'_L \neq \Gamma_L$. Nevertheless, the distortion can be made negligibly small. Indeed, using the quasi-classical Gamow formula, $\Gamma_L = (1/T_1)\exp(-2S)$, where $S = [2m(V - E_1)]^{1/2}L$ and $T_1 = [2m/E_1]^{1/2}L_1$, one finds

$$\frac{\Gamma'_L - \Gamma_L}{\Gamma_L} = \frac{U}{2E_1} + S\frac{U}{V - E_1}. \tag{6}$$

It follows that in the limit $U/E_1 \to 0$ and $U/(V - E_1) \to 0$, the width $\Gamma'_L \to \Gamma_L$. As a result $I_S \to I_S^{(0)}$, so that the measurement of I_D can be considered as a *noninvasive* measurement of the current I_S.

In general, the measurement is a noninvasive one if the density matrix of the measured object does not depend on the detector parameters. One can easily see that this is precisely

the case for the measurement discussed above. Indeed, by introducing the density matrix of the measured system, $\bar{\sigma}(t)$, obtained by tracing out the detector states in $\sigma(t)$, we find from Eq. (4) in the limit $\Gamma'_{L,R} \to \Gamma_{L,R}$

$$\dot{\bar{\sigma}}_{aa} = -\Gamma_L\bar{\sigma}_{aa} + \Gamma_R\bar{\sigma}_{cc} \tag{7a}$$

$$\dot{\bar{\sigma}}_{cc} = \Gamma_L\bar{\sigma}_{aa} - \Gamma_R\bar{\sigma}_{cc} \tag{7b}$$

where $\bar{\sigma}_{aa} = \sigma_{aa} + \sigma_{bb}$ and $\bar{\sigma}_{cc} = \sigma_{cc} + \sigma_{dd}$. The detector parameters Γ_0 and Γ'_0 are thus canceled in the equation of motion of the measured system. This shows that a noninvasive measurement can be performed even if the measured system is not a macroscopic object [2,3].

Note that the attachment of the detector to the measured system can be considered the last step in the measurement process. Even though one still need to register the current in the detector's right reservoir, this additional process can be carried out without any further distortion of the measured system and even the detector. Indeed, the reservoirs are systems with continuum spectrum, described by the classical rate equations that, as we have seen, admit a noninvasive measurement.

Consider now resonant transport in the coupled well (coupled dot) structure shown in Fig. 2. For simplicity, we assume that the Coulomb repulsion inside the double-dot is very strong, so only one electron can occupy it [12]. First consider the case of $E_0 + U_1$, $E_0 + U_2 > \tilde{E}_F^L$. Then an electron cannot enter the detector whenever another electron occupies the double-dot. All possible states of the measured system and the detector are shown in Fig. 2. We assume from the very beginning that the distortion of the coupled-well parameters by the detector is negligibly small, i.e., $\Gamma'_{L,R} \to \Gamma_{L,R}$, $\Omega' \to \Omega$. Using Eqs. (3) we can write the corresponding rate equations for the density matrix

$$\dot{\sigma}_{aa} = -(\Gamma_L + \Gamma_0)\sigma_{aa} + \Gamma_0\sigma_{bb} + \Gamma_R\sigma_{dd} \tag{8a}$$

$$\dot{\sigma}_{bb} = -(\Gamma_L + \Gamma_0)\sigma_{bb} + \Gamma_0\sigma_{aa} + \Gamma_R\sigma_{ff} \tag{8b}$$

$$\dot{\sigma}_{cc} = i\Omega(\sigma_{cd} - \sigma_{dc}) + \Gamma_L\sigma_{aa} + 2\Gamma'_0\sigma_{ee} \tag{8c}$$

$$\dot{\sigma}_{dd} = -i\Omega(\sigma_{cd} - \sigma_{dc}) - \Gamma_R\sigma_{dd} + 2\Gamma_0''\sigma_{ff} \tag{8d}$$

$$\dot{\sigma}_{ee} = i\Omega(\sigma_{ef} - \sigma_{fe}) - 2\Gamma_0'\sigma_{ee} + \Gamma_L\sigma_{bb} \tag{8e}$$

$$\dot{\sigma}_{ff} = -i\Omega(\sigma_{ef} - \sigma_{fe}) - (\Gamma_R + 2\Gamma_0'')\sigma_{ff} \tag{8f}$$

$$\dot{\sigma}_{cd} = i\epsilon\sigma_{cd} + i\Omega(\sigma_{cc} - \sigma_{dd}) - \frac{1}{2}\Gamma_R\sigma_{cd} + (\Gamma_0' + \Gamma_0'')\sigma_{ef} \tag{8g}$$

$$\dot{\sigma}_{ef} = i(\epsilon + \Delta U)\sigma_{ef} + i\Omega(\sigma_{ee} - \sigma_{ff}) - (\Gamma_0' + \Gamma_0'' + \Gamma_R/2)\sigma_{ef}, \tag{8h}$$

where $\epsilon = E_2 - E_1$ and $\Delta U = U_2 - U_1$. The currents in the detector (I_D) and in the double-dot system (I_S) are given by Eq. (2): $I_D/e = \Gamma_0\sigma_{bb} + \Gamma_0'\sigma_{ee} + \Gamma_0''\sigma_{ff}$, and $I_S/e = \Gamma_R(\sigma_{dd} + \sigma_{ff})$. Solving Eqs. (8) we find

$$I_S = I_S^{(0)}(1 - \alpha I_S^{(0)}/\Gamma_0), \qquad\qquad I_D/I_S = \Gamma_0/2\Gamma_L, \tag{9}$$

where

$$I_S^{(0)}/e = \frac{\Gamma_R\Omega^2}{\epsilon^2 + \Gamma_R^2/4 + \Omega^2(2 + \Gamma_R/\Gamma_L)} \tag{10}$$

is the undistorted resonant current in the double-dot system for a strong Coulomb repulsion [8,12]. Consider $\Gamma_0 \sim \Gamma_0' \sim \Gamma_0'' \gg \Gamma_{L,R}, \Omega, \epsilon$. In this limit the coefficient α in Eq. (9) is $\alpha = \Delta U(\Delta U + \epsilon)/(4\Gamma_0)^2$. As a result, $I_S \to I_S^{(0)}$ for $\Delta U \ll \Gamma_0$. Thus, the measurement of the detector current can be considered as a noninvasive measurement of the double-dot current. Also, one easily checks that the density matrix of the double-dot system, $\bar{\sigma}(t)$, is decoupled from the detector in the limit $\Gamma_{L,R}' \to \Gamma_{L,R}$, $\Omega' \to \Omega$ and $\Delta U \to 0$. Indeed, one finds from Eq. (8)

$$\dot{\bar{\sigma}}_{aa} = -\Gamma_L\bar{\sigma}_{aa} + \Gamma_R\bar{\sigma}_{dd} \tag{11a}$$

$$\dot{\bar{\sigma}}_{cc} = i\Omega(\bar{\sigma}_{cd} - \bar{\sigma}_{dc}) + \Gamma_L\bar{\sigma}_{aa} \tag{11b}$$

$$\dot{\bar{\sigma}}_{dd} = -i\Omega(\bar{\sigma}_{cd} - \bar{\sigma}_{dc}) - \Gamma_R\bar{\sigma}_{dd} \tag{11c}$$

$$\dot{\bar{\sigma}}_{cd} = i\epsilon\bar{\sigma}_{cd} - i\Omega(\bar{\sigma}_{cc} - \sigma dd) - \frac{1}{2}\Gamma_R\bar{\sigma}_{cd}, \tag{11d}$$

where $\bar{\sigma}_{aa} = \sigma_{aa} + \sigma_{bb}$, $\bar{\sigma}_{cc} = \sigma_{cc} + \sigma_{ee}$, $\bar{\sigma}_{dd} = \sigma_{dd} + \sigma_{ff}$ and $\bar{\sigma}_{cd} = \sigma_{cd} + \sigma_{ef}$.

The above example shows that the behavior of the measured system is not distorted by the measurement, even if the system is in a linear superposition of different states, which affect of the detector in a different way (compare the states c and d, or e and f in Fig. 2). Notice that the detector remains blocked, whenever an electron occupies any of the wells of the measured system. As a result such a measurement cannot single out the well in which an electron is located.

Let us now increase the Fermi energy $\tilde{E}_F^L$, so that $E_0 + U_2 < \tilde{E}_F^L < E_0 + U_1$. In this case an electron from the left reservoir can enter the detector even when the second dot of the measured system is occupied (the states (d) and (f) in Fig.2). As a result, the rate equations for $\sigma_{dd}, \sigma_{ff}, \sigma_{cd}, \sigma_{ef}$ are changed. One finds

$$\dot{\sigma}_{dd} = -i\Omega(\sigma_{cd} - \sigma_{dc}) - (\Gamma_R + \Gamma_0'')\sigma_{dd} + \Gamma_0''\sigma_{ff} \tag{12a}$$

$$\dot{\sigma}_{ff} = -i\Omega(\sigma_{ef} - \sigma_{fe}) - (\Gamma_R + \Gamma_0'')\sigma_{ff} + \Gamma_0''\sigma_{dd} \tag{12b}$$

$$\dot{\sigma}_{cd} = i\epsilon\sigma_{cd} + i\Omega(\sigma_{cc} - \sigma_{dd}) - \frac{1}{2}(\Gamma_R + \Gamma_0'')\sigma_{cd} + \frac{1}{2}(\Gamma_0' + \Gamma_0'')\sigma_{ef} \tag{12c}$$

$$\dot{\sigma}_{ef} = i(\epsilon + \Delta U)\sigma_{ef} + i\Omega(\sigma_{ee} - \sigma_{ff}) - \frac{1}{2}(\Gamma_R + \Gamma_0' + 2\Gamma_0'')\sigma_{ef}, \tag{12d}$$

instead of Eqs. (8d),(8f)-(8h). Such a modification does not affect the diagonal density matrix of the measured system, $\bar{\sigma}_{ii}(t)$, Eqs. (11a)-(11c). However, the rate equation for the non-diagonal density matrix element, $\bar{\sigma}_{cd}$, is different from Eq. (11d). We obtain from Eqs. (12c)-(12d) for $\Delta U = 0$

$$\dot{\bar{\sigma}}_{cd} = i\epsilon\bar{\sigma}_{cd} - i\Omega(\bar{\sigma}_{cc} - \sigma_{dd}) - \frac{1}{2}\Gamma_R\bar{\sigma}_{cd} - \frac{1}{2}(\Gamma_0''\bar{\sigma}_{cd} + \Gamma_0'\bar{\sigma}_{ef}) \tag{13}$$

Thus, the widths Γ_0', Γ_0'' are not canceled out in the density matrix of the double-dot system, despite the fact that no parameter of this system is distorted by the detector. Moreover, the influence of the detector on the measured current, I_S, is very strong. Indeed, take for simplicity $\Gamma_0'' \simeq \Gamma_0'$. Then $I_S = I_S^{(1)}$, where

$$I_S^{(1)}/e = \frac{\Gamma_R\Omega^2}{\epsilon^2\Gamma_R/(\Gamma_R + \Gamma_0') + \Gamma_R(\Gamma_R + \Gamma_0')/4 + \Omega^2(2 + \Gamma_R/\Gamma_L)} \tag{14}$$

It follows from Eq. (14) that $I_S^{(1)} \to 0$ for $\Gamma_0' \gg \Gamma_{L,R}$, Ω, ϵ. However, the current in the double-dot system returns to a previous (non-distorted) value, $I_S^{(0)}$, Eq. (10), for $\tilde{E}_F^L >$

$E_0 + U_1$, Fig. 3. (We predict the same behavior of the current I_S as a function of $\tilde{E}_F^L$ if the detector is located near the second dot).

The expected strong drop-off of the current I_S for $E_0 + U_1 > \tilde{E}_F^L > E_0 + U_2$ can be interpreted as an "observation" effect, yet without an "observer". Indeed, this configuration of the Fermi level allows us, in principle, to distinguish a particular dot occupied by an electron, whereas other configurations do not. It is clear that this effect cannot be explained by an interaction between electrons inside the detector and the measured system, since the distortion of the measured system due to this interaction is negligibly small. Even if the distortion of the system cannot be totally neglected, the same distortion would also exist for $\tilde{E}_F^L < E_0 + U_1$ or $\tilde{E}_F^L > E_0 + U_2$, where no effect is expected.

Actually, our quantum rate equations point to the origin of the observation effect. Consider first Eq. (3a), describing the diagonal density matrix elements. One finds that for each transition $a \to a'$ there exists the reverse $a' \to a$, which contributes with the opposite sign to the rate equation for $\sigma_{a'a'}$. Therefore the corresponding rates Γ cancel in the rate equation for $\sigma_{aa} + \sigma_{a'a'}$. As a result, the detector rates drop out from the density matrix when the latter is traced over the detector states (provided the rates of the measured system are not distorted by the measurement). However, the rates Γ are not in balance in the rate equation for non-diagonal density-matrix elements, Eq. (3b). One easily finds that the negative contributions from the transitions $(ab) \to (a'b)$ and $(ab) \to (ab')$ have no positive counterparts (in contrast with the transitions $(ab) \to (a'b')$) [6]. This is precisely the case of measurement of a system in a coherent superposition when one of the superposed states can generate or prevent transitions between continuum and isolated states in the detector while the other state does not. As a result, the corresponding detector rates are not canceled in the density matrix describing the measured system, even though the distortion of this system is negligibly small.

Although we concentrated in this paper on the measurement of currents in mesoscopic systems, we assert that the results are valid for the measurement problem in general. First,

consider the detector and the role of an observer. Since we found no dividing line between a microscopic and a macroscopic (classical) description, the detector need not be a macroscopic object. We require only the absence of transitions between discrete states, so that the diagonal and nondiagonal density matrix elements of the detector are decoupled. In this case the subsequent (noninvasive) interactions with an observer do not change the detector behavior, and therefore cannot influence the wave function collapse. It follows that the result of the measurement is not affected by the observer.

With respect to measurements of systems in a linear superposition, we found that the measurement process is fully reproduced by the Schrödinger equation, written as quantum rate equations for the density matrix, without an independent projection postulate [1]. Indeed, one easily obtains from Eqs. (11),(13) that $\bar{\sigma}_{cd} \to 0$ for $\Gamma'_0 \to \infty$, so that the density matrix of the double-dot system collapses into the corresponding statistical mixture. Notice, however, that for a finite Γ'_0, the nondiagonal density matrix elements $\bar{\sigma}_{cd}, \bar{\sigma}_{dc}$ do not disappear, but are only diminished. Therefore the damping of nondiagonal density matrix elements due to the environment does not necessarily lead to their elimination for $t \to \infty$, as suggested in [13].

Although the measurement process is described by the Schrödinger equation, the observation paradox is still there. It appears that a very sensitive detector, far away from the measured system but still capable of observing different states from the linear superposition, would strongly affect the measured system. This situation resembles the EPR paradox, but some features are different. First, the above observation paradox appears as a stationary state phenomenon. Second, we do not need any special initial correlations between the electrons in the detector and the measured system. The most important difference is the possibility of influencing directly the measured current by switching the detector on (or off). Such a process can also be studied using our rate equations (3), describing time dependence of the density matrix, but special attention should be paid to the relativistic treatment [14].

ACKNOWLEDGMENTS

I am grateful to Ya. Prager and B. Svetitsky for reading the manuscript and making valuable comments on it.

REFERENCES

[1] J. von Neumann, *Mathematische Grundlagen der Quantentheorie* (Springer, Berlin, 1931). The statistical mixture means that the measured system is actually in one of its states, although there is no way to determine in which particular state the system is. Otherwise, quantum mechanics would be a pure deterministic theory.

[2] A.J. Leggett and A. Garg, Phys. Rev. Lett. **54**, 857 (1985); Phys. Rev. Lett. **59**, 1621 (1987); Phys. Rev. Lett. **63**, 2159 (1989). The predicted violation of the temporal Bell inequalities has not yet been observed.

[3] For a realization of interaction-free measurement see L. Vaidman, Quant. Opt. **6**, 119 (1994); P. Kwiat, H. Wenfurter, T. Herzog, A. Zeilinger and M. Kasevich, Phys. Rev. Lett. **74**, 4763 (1995).

[4] S.A. Gurvitz, H.J. Lipkin, and Ya.S. Prager, Mod. Phys. Lett. B **8**, 1377 (1994).

[5] S.A. Gurvitz, H.J. Lipkin, and Ya.S. Prager, Phys. Lett. A, in press.

[6] S.A. Gurvitz and Ya.S. Prager, cond-mat/9511026, Phys. Rev. B, in press.

[7] D.V. Averin and K.K. Likharev, in *Mesoscopic Phenomena in Solids*, edited by B. Altshuler, P.A. Lee, and R.A. Webb (Elsevier, Amsterdam, 1991).

[8] N.C. van der Vaart, S.F. Godijn, Y.V. Nazarov, C.J.P.M. Hartmans, J.E. Mooij, L.W. Molenkamp, and C.T. Foxon, Phys. Rev. Lett. **74**, 4702 (1995).

[9] F.R. Waugh, F.R. Waugh, M.J. Berry, D.J. Mar, R.M. Westervelt, K.L. Campman and A.C. Gossard, Phys. Rev. Lett. **75**, 4702 (1995).

[10] M. Field, C.G. Smith, M. Pepper, D.A. Ritchie, J.E.F. Frost, G.A.C. Jones, and D.G. Hasko, Phys. Rev. Lett. **70**, 1311 (1993).

[11] L.W. Molenkamp, K. Flensberg and M. Kemerink, Phys. Rev. Lett. **75**, 4282 (1995).

[12] Yu.V. Nazarov, Physica B **189**, 57 (1993), T.H. Stoof and Yu.V. Nazarov, Phys. Rev. B **55**, 1050 (1996).

[13] W.H. Zurek, Physics Today **44**, No. 10, 36 (1991); *ibid*, **46**, No. 4, 13 (1993), and references therein.

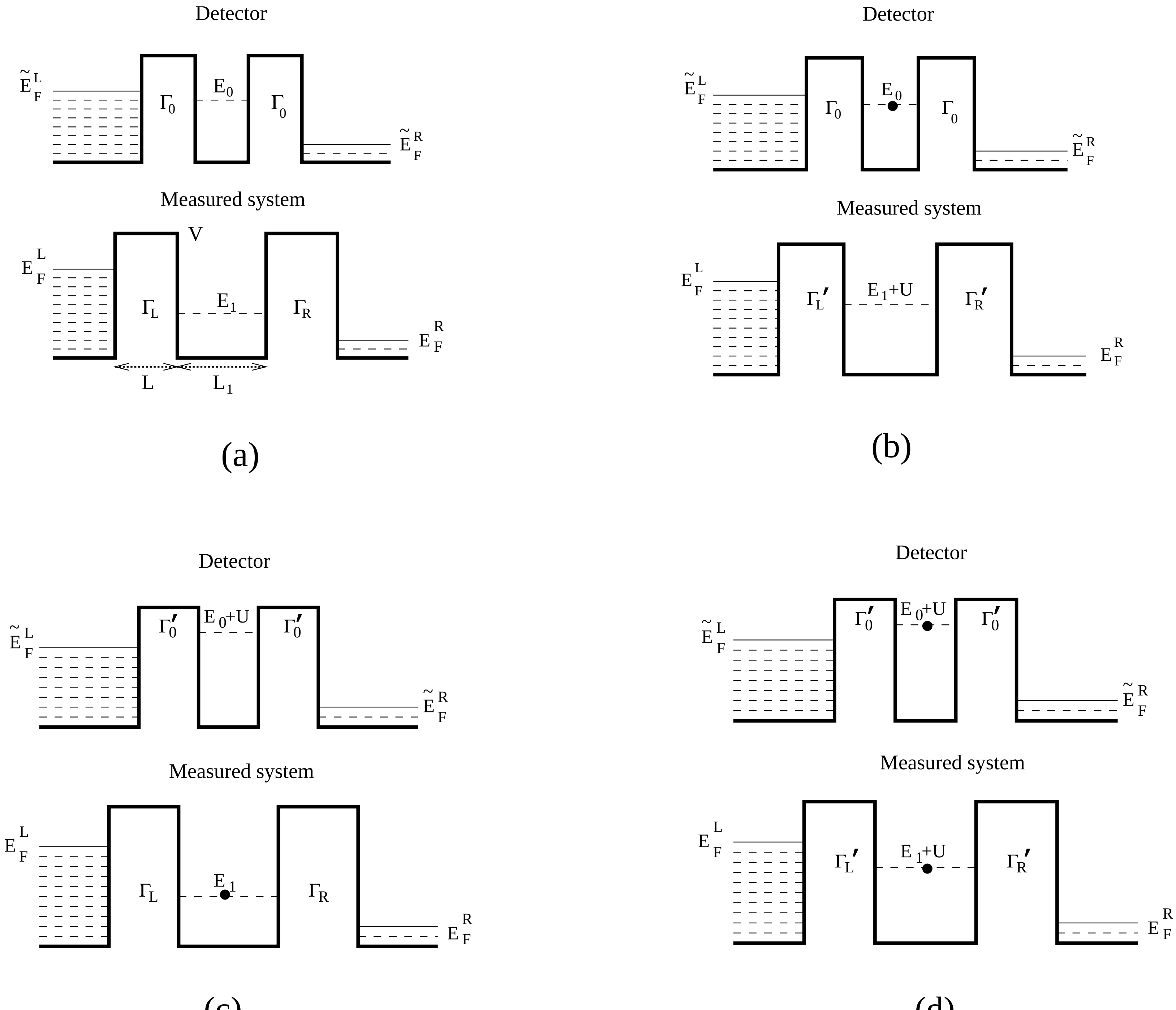
[14] Y. Aharonov, D. Albert and L. Vaidman, Phys. Rev. D **34**, 1805 (1986).

FIGURES

FIG. 1. The measurement of resonant current in a single-dot structure by another, nearby dot. All possible electron states of the detector (the upper well) and the measured system (the lower well) are shown. Also indicated are the tunneling rates (Γ), the left barrier height (V), barrier width (L), and width (L_1) of the lower well.

FIG. 2. The measurement of resonant current in a double-dot structure. The energy level of the upper dot (the detector) is above the Fermi level whenever an electron occupies one of the dots in the double-dot structure. Ω is the coupling between the dots.

FIG. 3. Maximal current in the double dot structure $I_S^{max} = I_S(\epsilon = 0)$ as a function of the Fermi energy of the left reservoir adjacent to the detector. The detector does not distort any parameters of the measured system.



Detector
Measured system
V
L
L_1
Γ_0
E_0
E_1
E_0+U
E_1+U
Γ_L
Γ_R
Γ_0'
Γ_L'
Γ_R'
$\tilde{E}_F^L$
$\tilde{E}_F^R$
E_F^L
E_F^R
(a)
(b)
(c)
(d)

Fig. 1

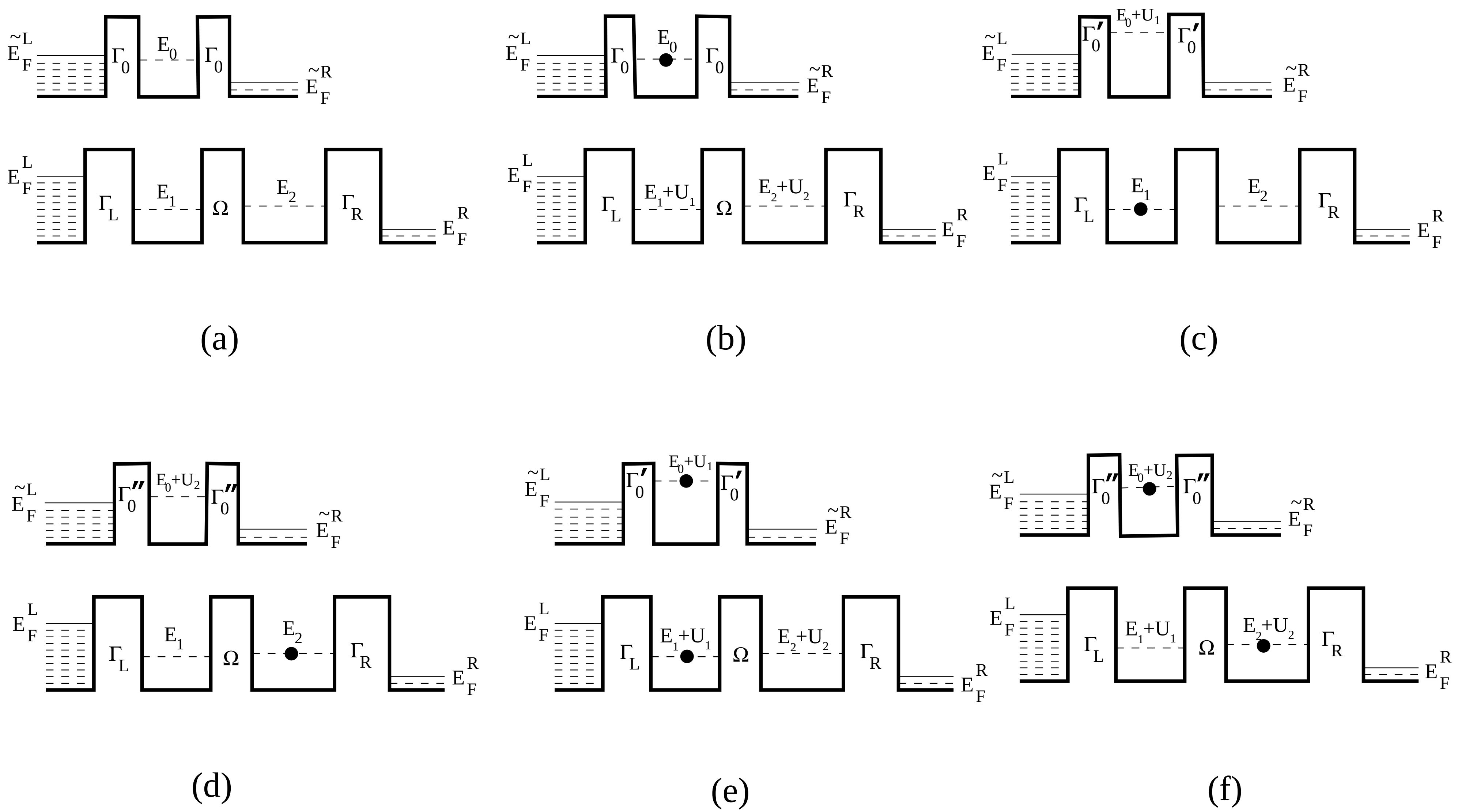


Fig. 2

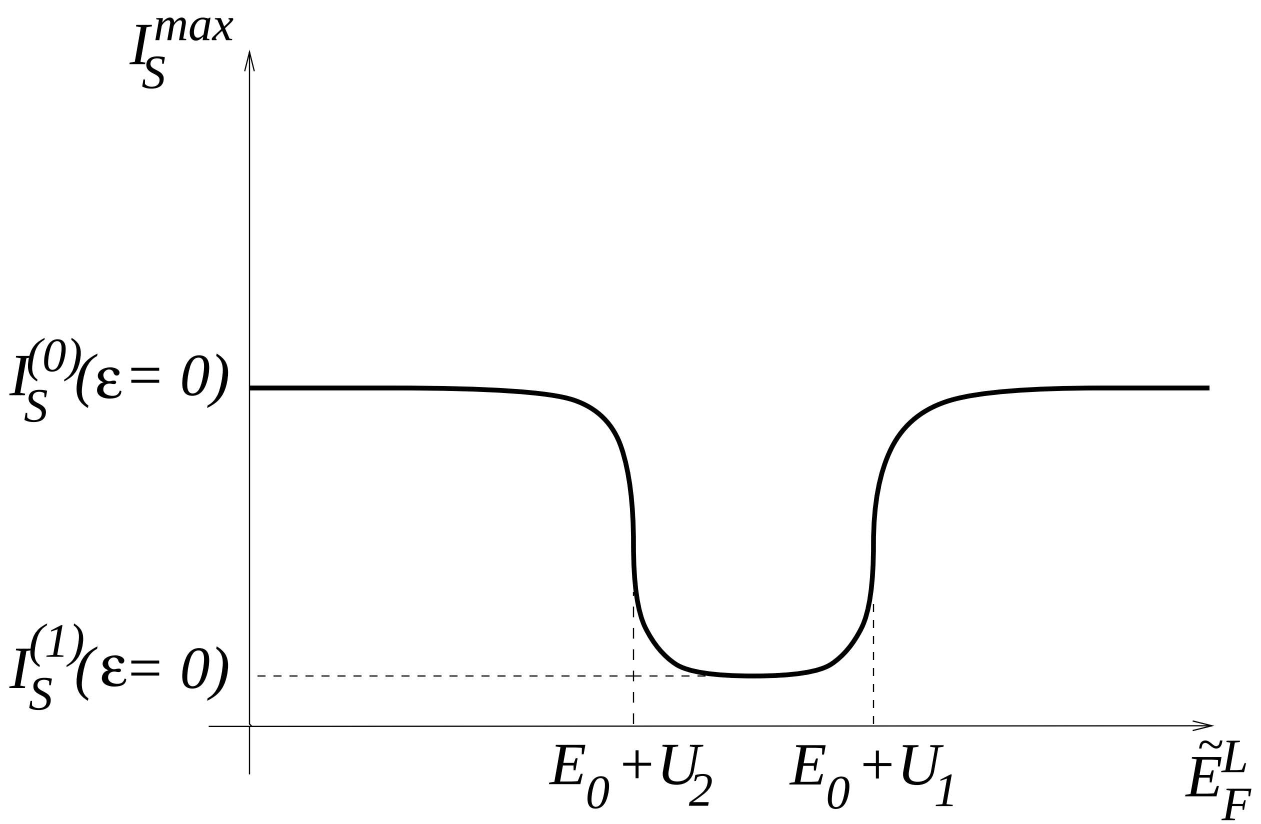


Fig. 3